

“Black Universe” epoch in String Cosmology

Alex Buchel^{1,2} and Lev Kofman³

¹*Department of Applied Mathematics*

University of Western Ontario

London, Ontario N6A 5B7, Canada

²*Perimeter Institute for Theoretical Physics*

Waterloo, Ontario N2J 2W9, Canada

³*CITA, University of Toronto*

60 St. George st.

Totonto, Ontario M5S 3H8, Canada

Abstract

String theory compactification involves manifolds with multiple warp factors. For cosmological applications, we often introduce a short, high-energy inflationary throat, and a long, low-energy Standard Model throat. It is assumed that at the end of inflation, the excited Kaluza-Klein modes from the Inflationary throat tunnel to the SM throat and reheat Standard Model degrees of freedom, which are attached to probe brane(s). However, the huge hierarchy of energy scales can result in a highly dynamic transition of the throat geometry. We point out that in such a cosmological scenario the Standard Model throat (together with SM brane) will be cloaked by a Schwarzschild horizon, produced by the Kaluza-Klein modes tunneling from the short throat. The Black Brane formation is dual to the first order chiral phase transition of the cascading gauge theory. We calculate the critical energy density corresponding the formation of the the BH horizon in the long throat. We discuss the duality between “Black Universe” cosmology and an expanding universe driven by the hot gauge theory radiation. We address the new problem of the hierarchical multiple-throat scenarios: SM brane disappearance after the decay of the BH horizon.

April 3, 2008

1 Introduction: Warped geometry, Duality and Cosmology

There are various high dimensional formulations of particle physics theories, such as fundamental string theory or phenomenological high dimensional constructions. Yet, most often string theory cosmology is reduced to the conventional $3 + 1$ dimensional universe, perhaps with some unconventional ingredients. Low energy effective description is obtained from the high dimensional theory by integration over compact inner dimensions

$$S_4 = \int d^4x \sqrt{-g} \mathcal{L}_4 = \int d^4x \sqrt{-g} \int d^6y \sqrt{G} \mathcal{L}_{10} . \quad (1.1)$$

An example recipe for string theory cosmology is to start with KKLT [1] (or large volume) compactification in the type IIB theory, include multiple throats with conifolds (attached to the bulk Calabi-Yau) to provide the hierarchy of masses and couplings, and to engineer inflation in this $4 + 6$ background geometry with branes or moduli fields. In such a cosmological scenario at the end point of inflation, reheating, energy is supposed to stream through the labyrinth of the compact manifold to find, in one of its corners, the light standard model particles which eventually heat and fill up the universe [3, 5, 6, 7]. For instance, one of the scenario is a warped brane inflation which occurs in the short (inflationary) throat at the GUT energy scale which contains an anti-brane $\overline{D3}$ at the tip attracting the mobile brane $D3$ [9]. Brane-antibrane annihilation terminates inflation and energy cascades into KK modes associated with the short throat. Suppose the SM particles are localized at the probe brane(s) in another, long SM throat of TeV scale. It is often assumed that KK modes of the short throat tunnel into the long throat and subsequently transfer their energy into SM sector. All the calculations can be done in ten dimensional warped throat geometry (with the radial coordinate y along of the throat)

$$ds^2 = H(y)^{-1/2} (-dt^2 + d\vec{x}^2) + H(y)^{1/2} G_{ab} dy^a dy^b , \quad (1.2)$$

in the supergravity approximation and then reduced to the effective four dimensional picture, according to the prescription (1.1). For the warped throats, one can use well-studied Klebanov-Strassler [10] solution with the deformed conifold at the tip, to estimate KK masses, tunneling and decay rates, etc.

The ten dimensional warped throat supergravity solution (like KS geometry) has a dual four dimensional gravity-free gauge theory description. Two or more throats

attached to the bulk CY will be dual to two or more gauge theories

$$\mathcal{L}(\phi_1, \phi_2) = \mathcal{L}_1(\phi_1) + \mathcal{L}_2(\phi_2) + \Delta\mathcal{L}(\phi_1, \phi_2), \quad (1.3)$$

where $\Delta\mathcal{L}(\phi_1, \phi_2) = \sum_n \frac{\mathcal{O}_n(\phi_1, \phi_2)}{M^n}$ is an interaction term due to the high dimensional operators, ϕ_1, ϕ_2 are the fields in two gauge theories. Excitations of KK modes in supergravity correspond to the glueball excitations in the dual picture (1.3), while the tunneling of KK modes between the throats has a dual description in terms of the high dimensional operator $\Delta\mathcal{L}(\phi_1, \phi_2)$.

A specific feature of the cosmological onset is that the two throats containing interacting excitations is a highly dynamic system with the huge hierarchy of the energy scales. Indeed, each of the gauge theories $\mathcal{L}_1(\phi_1), \mathcal{L}_2(\phi_2)$ has its own critical energy ϵ_c (or critical temperature T_c) which separates the lower energy confinement phase from the higher energy deconfinement plasma phase. On the other hand, the phase transitions at ϵ_c correspond to the emergence of the strong gravity regime in the supergravity dual, manifested in the appearance of a black hole (black brane) horizon across the throat, heuristically similar to the AdS/BH solution

$$ds^2 = \frac{r^2}{R^2} \left(- \left(1 - \frac{r_g^4}{r^4} \right) dt^2 + d\vec{x}^2 \right) + R^2 \frac{dr^2}{r^2 (1 - r_g^4/r^4)} + R^2 d\Omega_5^2. \quad (1.4)$$

where r is related to the radial direction y in (1.2). Therefore, to understand string cosmology scenario with two throats of the inflationary and SM scales one has to go well beyond the simple picture of KK modes tunneling between throats and to recall the higher dimensional description of the throat geometry which can contain the BH horizon. Moreover, the time-dependent cosmological scenario addresses to the string theory the questions how the BH horizon in the throat appears, evolves and eventually disappears, and what may be the interesting consequences of this higher dimensional epoch of the evolution of the universe. This is very different from the simple picture of KK modes tunneling between throats. Moreover, as the BH horizon screens the tip of the SM throat (perhaps also the $D3$ branes that live in the throat), the dy integration in Eq. (1.1) is not valid so that we have to deal with the higher dimensional theory. These are the problems we will try to formulate and discuss in this paper.

Let us recall the story of the high temperature phase transition /BH horizon duality. Although the motivation of each original study of such a duality was different, we will look at those from the perspective of our string theory cosmology scenario.

Phase transition/BH duality is deeply connected with the correspondence between supergravity in the $AdS_5 \times S^5$ and CFT at high temperature [11]. In the context of supergravity throat solutions and gauge theories, gravity dual to the supersymmetric $SU(K) \times SU(K + M)$ gauge theory¹ is the Klebanov-Tseytlin throat solution in the type IIB theory with fluxes [12] which has a singularity at the conifold. In [13] it was proposed that gauge theory with the restored chiral symmetry above the critical temperature T_c corresponds to the Schwarzschild horizon which cloaks the naked singularity of the KT solution or tip of the throat in the KS solution. While the original purpose in [13] was to resolve the singularity of the KT solution, its implications may go well beyond this. Klebanov-Strassler solution with deformed conifold resolves the singularity problem of KT solution at zero temperature. However, it is the BH horizon that provides a resolution of the KT singularity at high temperature. Technically, it is a challenge to find analytic BH throat solutions in ten dimensional type IIB theory with fluxes. Supergravity solution with the regular BH horizon in the KT throat was constructed in [14, 15, 16]. Gauge theory with the broken chiral symmetry at low temperature $T \ll T_c$ is dual to the regular throat solution, while gauge theory with restored chiral symmetry at $T \gg T_c$ is dual to the generalized throat solution with the regular BH horizon across the throat. Finite temperature gauge fields dynamics corresponds to the first order phase transition at T_c .

Originally, the KS throat solution was constructed as a supergravity dual to the cascading $SU(K) \times SU(K + M)$ gauge theory. However, because this self-consistent geometry describes the throat with the regular tip at the finite distance, and because this throat can be smoothly embedded in the CY manifold [17, 18], KS solution is often evoked by cosmologists as the model of the throat attached to the compact space.

Correspondence between BH horizon and the phase transition was also discussed in the very different context, namely, in the phenomenological five-dimensional Randall-Sundrum model. While usually the four-dimensional effective theory at the brane is derived from the bulk+branes system in the spirit of Eq. (1.1), there was a proposal in [19] to extrapolate the ideas of holography for the RS braneworld and particle phenomenology on the brane. Original high-temperature plasma phase was proposed to be dual to the bulk AdS/BH horizon in five dimensions, while EW phase transition at the brane should be dual to the disappearance of the horizon [20].

In the following we recall the cascading gauge theory side of the duality. After that

¹Assuming unbroken chiral symmetry.

we explicitly evaluate the critical energy density ϵ_c for a given throat. This will allow us to lay down the new cosmological scenario which contains higher dimensional stage of formation and later disappearance of the higher dimensional BH horizon. Similarly, this dynamics is relevant for the theoretical “thought experiment”, if one is not interested in cosmology. Finally we discuss potential cosmological problems and consequences of this cosmological model and address interesting theoretical issues.

2 Gauge Theory Perspective

The gauge theory dual to a local warped geometry is a so called “cascading gauge theory” introduced in [21, 12, 10]. This cascading gauge theory can be thought of as a specific $SU(K) \times SU(K + M)$ $\mathcal{N} = 1$ supersymmetric gauge theory, with a number of colors K which runs logarithmically with the energy scale μ [10, 13, 22, 23]

$$K = K(\mu) \sim M^2 \ln(\mu/\Lambda) , \quad (2.1)$$

where Λ is the strong coupling scale of the cascading gauge theory. Despite having an infinite number of degrees of freedom in the ultraviolet, cascading gauge theory is holographically² renormalizable as a four dimensional quantum field theory [24]. At a given scale μ the gauge theory has two chiral superfields A_1, A_2 in the $(K + M, \overline{K})$ representation, and two fields B_1, B_2 in the $(\overline{K + M}, K)$ representation. The superpotential of the model is

$$W \sim \text{Tr} (A_i B_j A_k B_\ell) \epsilon^{ik} \epsilon^{j\ell} . \quad (2.2)$$

The two gauge group factors have gauge couplings g_1 and g_2 . Under the renormalization group flow the sum of the coupling does not run

$$\frac{4\pi}{g_1^2} + \frac{4\pi}{g_2^2} = \text{constant} , \quad (2.3)$$

while the difference is

$$\frac{4\pi}{g_2^2} - \frac{4\pi}{g_1^2} \sim M \ln(\mu/\Lambda) [3 + 2(1 - \gamma)] , \quad (2.4)$$

where γ is the anomalous dimension of operators $\text{Tr} A_i B_j$. It is clear from (2.3), (2.4) that starting at some energy scale and flowing either to the UV or the IR one inevitably

²In the UV the cascading gauge theory ’t Hooft coupling becomes strong and thus renormalization of this gauge theory must be addressed in the framework of the dual gravitational description.

encounters a Landau pole: one of the two gauge couplings will become infinitely large. In [10] it was argued that extension of the RG flow past the infinite couplings is achieved by a cascade of self-similar Seiberg duality [25] transformations on the strongly coupled gauge group factor. At each duality step $K \rightarrow K + M$ for the RG flow to the UV, and $K \rightarrow K - M$ for the RG flow to the IR, leading to the effective logarithmic running of the number of colors (2.1).

Cascading gauge theory confines in the IR; it has a classical $U(1)_R$ symmetry which is explicitly broken to \mathbf{Z}_{2M} by the anomaly, and then it is broken spontaneously by the gluino condensate to \mathbf{Z}_2 [10]. In [13] it was pointed out that at sufficiently high temperature, the R-symmetry of the gauge theory is restored; moreover, this restoration is accompanied by a first order deconfinement phase transition³. Such a phase transition was recently identified in [16]. A summary of the proposal [13] and the detailed analysis [16] is that a cascading gauge theory at equilibrium has a critical energy density, which we refer to as ϵ_c , such that for density $\epsilon > \epsilon_c$ it is in a deconfined phase with $\mathcal{O}(K(\epsilon)^2)$ entropy.

Several comments are in order before we finish the gauge theory discussion.

First, the deconfinement phase transition at $\epsilon > \epsilon_c$ assumes that the boundary gauge theory is in flat space-time, or at least $\epsilon_c \gg R_4^2$. Here, R_4 is the Ricci scalar of the background metric.

One expects a chiral symmetry restoration (but *without* the deconfinement) phase transition in curved space-time. For a cascading gauge theory on S^3 this was demonstrated in [26] while for the cascading gauge theory in dS_4 this was discussed in [27]. In both cases the background curvature serves as a regulator that cuts off the IR physics associated with the gluino condensate. By studying a $D3$ -brane probe in the deformed geometry⁴ it is easy to see that the spectrum of fluctuations in the deformed throat is gapped.

In principle, for a cascading gauge theory in flat space-time one could imagine two separate phase transitions: a confinement/deconfinement one, and the chiral symmetry restoration transition. For the scenario discussed in this paper it is important whether or not there is a deconfined phase with a broken chiral symmetry. Such a phase, were it to exist, is expected to have a critical energy density higher than ϵ_c . A detailed analysis

³The high temperature deconfinement state of the cascading gauge theory was studied in [13, 14, 15].

⁴In the dS_4 case this was done in [28].

[29] indicates that such a phase is not realized. In other words: the hot cascading gauge theory plasma cools through a first order phase transition where the chiral symmetry is broken only when the theory confines.

3 Critical Energy Density in the Throat

In this Section we turn to the throat geometry. We will compute the critical energy density for weakly curved 3+1 dimensional space-time (modeling our universe) that is necessary to hide it behind the horizon in the ambient warped throat geometry. Each throat has its own critical density, which corresponds to horizon formation, but for our scenario we will focus on ϵ_c for the SM throat.

We assume that our universe (in the low energy limit where the energy density in all throats is much lower than ϵ_c) is part of the KS geometry which away from the tip is locally described by the KT geometry

$$ds_{10}^2 = H^{-1/2}(r) \eta_{\mu\nu} dx^\mu dx^\nu + H^{1/2}(r) (dr^2 + r^2 ds_{T^{1,1}}^2) , \quad (3.1)$$

where

$$H(r) = \frac{R_+^4 + R_-^4 \ln\left(\frac{r}{R_+}\right)}{r^4} , \quad (3.2)$$

with

$$R_+^4 = \frac{27\pi}{4} \alpha'^2 g_s M K , \quad R_-^4 = \frac{81}{8} \alpha'^2 g_s^2 M^2 , \quad (3.3)$$

and M and K are the integer numbers associated with the quantization conditions of the form fields fluxes of the type IIB theory which generate the throat geometry (3.1). The local geometry (3.1) is holographically dual to the four-dimensional cascading gauge theory reviewed in the previous section. For large enough energy density this cascading gauge theory undergoes a deconfining phase transition [13, 16], which is reflected in the formation of a Schwarzschild horizon⁵ in (3.1). The type IIB supergravity background corresponding to a Schwarzschild horizon in asymptotic geometry

⁵As we will explain in the following Sections, in our scenario this energy density is accumulated from the “gravitational collapse” of the excited KK modes initially produced from the $D3\overline{D3}$ annihilation in the inflationary throat.

(3.1) takes form [16]⁶:

$$ds_{10}^2 = h^{-1/2} (2Y - Y^2)^{-1/2} \left(-(1 - Y)^2 dt^2 + d\vec{x}^2 \right) + G_{YY} dY^2 + h^{1/2} [f_2 (e_\psi^2) + f_3 \sum_{a=1}^2 (e_{\theta_a}^2 + e_{\phi_a}^2)], \quad (3.4)$$

where h , f_2 and f_3 are some functions of the radial coordinate $Y \in [0, 1]$. There is also a dilaton, and form fields (see [16] for details). Notice that the radial coordinate in (3.4) is gauged-fixed⁷ so that

$$\frac{G_{tt}}{G_{ii}} = -(1 - Y)^2. \quad (3.5)$$

The new radial coordinate Y is related asymptotically to r in (3.1) as follows

$$Y \sim \frac{1}{r^4}, \quad r \rightarrow \infty. \quad (3.6)$$

Type IIB supergravity equations of motion in the background metric (3.4) are solved, subject to the following boundary conditions:

- i) near the boundary ($Y \rightarrow 0$ or $r \rightarrow \infty$) the black hole metric (3.4) approaches the Klebanov-Tseytlin geometry (3.1);
- ii) the hypersurface $Y = 1$ is a regular Schwarzschild horizon of the metric (3.4); the latter is equivalent to requiring that all the warp factors h , f_2 and f_3 are positive at $Y = 1$.

When the temperature T of the black hole (3.4) is much larger than the characteristic KK scale m_{KK} deep inside the throat geometry (3.4)⁸

$$T \gg \Lambda \sim m_{KK} \sim \frac{1}{R_-} e^{-\frac{2\pi K}{3Mg_s}}, \quad (3.7)$$

the black hole geometry and its thermodynamics can be determined analytically [15, 24, 16]. Specifically, we find for $T \gg \Lambda$

$$f \simeq -\frac{1}{8} \pi^2 K(T)^2 T^4 \sim -M^4 T^4 \left(\ln \frac{T}{\Lambda} \right)^2, \quad \epsilon \simeq -3f, \quad (3.8)$$

⁶The frames $\{e_{\theta_a}, e_{\phi_a}\}$ are defined as in [24], such that the metric on a unit size $T^{1,1}$ is given by $(e_\psi^2) + \sum_{a=1}^2 (e_{\theta_a}^2 + e_{\phi_a}^2)$.

⁷Because we gauge-fixed the radial coordinate, there is a constraint equation coming from the equation of motion of this variable; this equation can be solved to determine G_{xx} [16].

⁸See Eq.(4.1) of the next Section.

for the black hole free energy density f and the energy density ϵ . We have used (2.1) to arrive at (3.8). Here $K(T)$ is the number of degrees of freedom, c.f. Eq. (2.1).

For temperature T of order Λ the black hole thermodynamics can be studied only numerically. It was found in [16] that there is a critical temperature, corresponding to a critical energy density ϵ_c , such that the free energy density of the black hole (3.4) vanishes precisely at $\epsilon = \epsilon_c$, and becomes positive for $\epsilon < \epsilon_c$. Thus we expect a first order deconfinement/confinement phase transition in the cascading gauge theory plasma once its energy density becomes less than ϵ_c . This first order phase transition (which occurs via bubble nucleation and percolation) shall be dual to the inhomogeneous melting of the BH horizon (which before this is translationally invariant in the \vec{x} directions).

From the point of view of the usual 3+1 dimensional GR, the disappearance of the horizon is rather unusual. However, the disappearance of a higher dimensional horizon is a familiar phenomenon, and we shall comment on it. A crucial observation is that at the supergravity side, there is another geometry, besides the black hole geometry (3.4), which also asymptotes to (a Euclidean version) of (3.1): the KS geometry with a Euclidean time direction identified with a period of $1/T$. The latter solution has no horizon, it has a zero free energy, and is dual to a thermal gas of the confined cascading gauge theory. The coexistence of these two solutions provides the possibility for transitions between horizon and no-horizon geometrical phases (through instantons, see.e.g. [20]). The supergravity dual of the phase transition bears similarity to the Hawking-Page transition, which was identified by Witten as a supergravity dual to a (kinematic) confinement/deconfinement transition of the $\mathcal{N} = 4$ SYM plasma on a three-sphere [11]. There is a notable difference however: while the Hawking-Page-Witten transition occurs in a finite volume⁹, the phase transition in the cascading gauge theory occurs in the infinite volume. Thus, we expect it to proceed via the nucleation of bubbles of the stable phase, which further expand and 'remove' the horizon.

Next, we evaluate ϵ_c quantitatively. In the notation of [16], a Schwarzschild horizon is formed in (3.1) at the critical energy density (at the gauge theory side) which is given by equation (see also Eq. (4.10) of [16])

$$\epsilon_c = \frac{1}{4\pi G_5} a_0^2, \quad (3.9)$$

⁹The transition disappears in the infinite volume.

where the 5-dimensional Newton's constant G_5 is

$$\frac{1}{G_5} = \frac{vol_{T^{1,1}}}{G_{10}} = \frac{16\pi^3}{27} \times \frac{1}{8\pi^6 g_s^2 \alpha'^4}, \quad (3.10)$$

and a_0 will be evaluated momentarily. A careful matching of the asymptotic black hole geometry in [16] with (3.1) leads to identification

$$\left(\frac{R_+}{R_-}\right)^4 - \frac{1}{4} \ln \frac{2eR_+^4}{a_0^2} = \frac{1}{2} k_c, \quad (3.11)$$

where

$$k_c = 0.25712(1), \quad (3.12)$$

(see Eq. (5.12) of [16]).

From (3.3)-(3.11) we conclude

$$\epsilon_c = 2MK e^{1+2k_c} T_3 e^{-4A}, \quad (3.13)$$

where T_3 is a $D3$ brane tension $T_3 = \frac{1}{(2\pi)^3 g_s \alpha'^2}$ and $e^{-A} = e^{-\frac{2\pi K}{3Mg_s}}$ is the hierarchy warp factor of the throat [30]. The four-dimensional energy-density of KK modes (located at the tip of the throat) is identified with the energy-density of the glueballs in the dual picture. Also notice that the bigger energy ϵ corresponds to the bigger portion of the throat cloaked by the horizon.

4 Theory of Multiple-throat Tunneling Revisited

Suppose compact manifold contains throats of significantly different warpings which are attached to the bulk, as sketched in the Figure 1. Suppose the short throat is associated with the higher energy scale and is populated with Kaluza-Klein excitations. It is well known that the typical KK masses are of the order of

$$m_{KK} \simeq \frac{e^{-A}}{R}, \quad (4.1)$$

where $e^{-A} = H(y)^{-1/4}$ in the warping factor at the tip of the throat, R is the radius scale of the angular coordinates at the tip, for the metric (3.1) it is R_- . Wave functions of the KK modes in the warped geometry are exponentially peaked around the tip of the throat, as illustrated in the Figure 1.

Let us set up the thought experiment with the static geometry of the Figure 1, with the flat outer space. Suppose the short throat is filled up with the KK excitations of

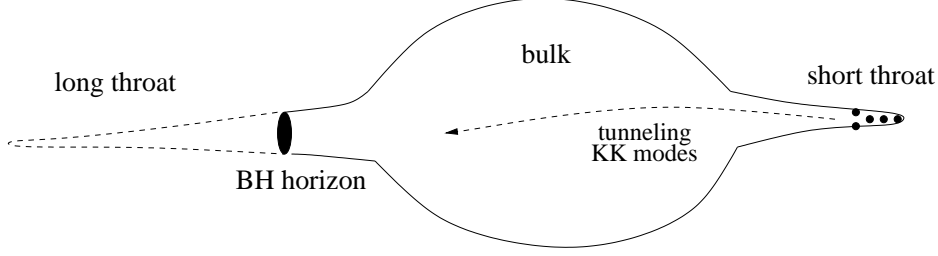


Figure 1: Sketch of the compact manifold with two throats.

the mass (4.1) and four-dimensional energy density ϵ_1 . For definiteness we can take $e^{-A_1} = 10^{-3}$, $R_1 \simeq 10\sqrt{\alpha'} \simeq 10^2 M_p^{-1}$ so that $m_{1KK} \sim 10^{13}$ GeV (index 1 attributes parameters to the short throat, while index 2 to the long throat). Suppose $\epsilon_1 \sim T_3 \sim (10^{16} \text{ GeV})^4$. KK modes from the short throat begin to tunnel into the long throat. Putting aside interesting subtleties (see e.g. [3, 6, 7, 8]) let us use the estimation of [31] for the inter-throat tunneling rate

$$\Gamma_{tun} \simeq (m_{1KK} R_1)^4 e^{-A_1} \sim \frac{e^{-5A_1}}{R_1}, \quad (4.2)$$

which gives us $\tau_{tun} = 1/\Gamma_{tun} \sim 10^{-26}$ sec. During this time the long throat is filling up with KK excitations, each with mass much lighter than m_{1KK} . If we choose $e^{-A_2} = 10^{-15}$, $R_2 = R_1$, then $m_{2KK} \sim 10$ GeV. Those modes again are accumulating¹⁰ around the tip of the long throat, with ever increasing four-dimensional energy density $\epsilon_2(t)$. However, gravitational backreaction of KK modes on the geometry of the long-throat becomes significant as soon as $\epsilon_2(t)$ is approaching from below the critical density ϵ_{2c} of the long throat. Let us estimate ϵ_{2c} from Eq. (3.13). For example, choosing $MK \sim 10^4$, from $\epsilon_{2c} \sim ((MK)^{1/4} 10 \text{ GeV})^4$ we get $\epsilon_{2c} \sim (100 \text{ GeV})^4$, i.e. the scale of the EW phase transition. This is much lower than the original energy density ϵ_1 of the excitations in the system. It means that long before the time τ_{tun} KK modes accumulating in the long throat will completely change the geometry of that throat. Actually, it happens instantly¹¹.

While it is immensely difficult to follow the detailed complicated metamorphose of the compact space in the self-consistent supergravity formalism, the dual picture in the gauge theory side suggests that increasing energy of the plasma above ϵ_{2c} corresponds

¹⁰We can neglect the tunneling of KK modes from the long (SM) throat back to the short (inflationary) throat as the latter is suppressed compare to (4.2) by a factor $\propto e^{5(A_2 - A_1)} \propto 10^{60}$.

¹¹This follows from the energy balance equation $\epsilon_2 = \epsilon_1 (1 - e^{-t\Gamma_{tun}})$.

to the formation of the BH horizon in the throat geometry. Since the chiral phase transition in the gauge theory plasma is a first order one, the formation of the BH horizon will be dual to the formation of the bubbles of the new phases which eventually percolate.

Therefore we conjecture the following picture of the throat BH formation. The horizon begins to cloak the tip of the throat and propagates, further screening a bigger and bigger part of the throat. This may occur in the manner of the Choptuik critical collapse [32]. Formation of the horizon may occur not uniformly in space, but in patches which later percolate into the uniform horizon (translationary-invariant in three-dimensional outer space). As KK modes from the short throat continue to tunnel into the long throat, BH there absorbs all of them and it happens much faster than the tunneling time $1/\Gamma_{tun}$ (this timing is defined by the cross-section of the high dimensional black hole).

All together, in our thought experiment we start with the excitations of the KK modes in the short throat and end up with the high dimensional Black Hole cloaking the tip of the long throat. From the four dimensional perspective, there are no particles, but uniformly distributed energy density from the high dimensional BH. Again, the dual picture suggest this form of the energy density has the radiation dominated equation of state but without radiation! Recall similar situation in the RS braneworlds with the AdS/Schwarzschild bulk geometry where four-dimensional “dark radiation” is associated with the projection of the five-dimensional Weyl tensor [33].

5 “Black Universe” Cosmological Scenario

The setting of the thought experiment of the previous Section takes place naturally in the popular string theory cosmological model. Indeed, in this model, the warped brane inflation is based on the brane-antibrane interaction in the short (inflationary) throat, which provides very shallow effective four-dimensional inflationary potential [9]. Four dimensional energy density of the branes is $\epsilon_1 = 2T_3$. At the end of inflation the brane-antibrane pair annihilates and releases energy into KK modes excitations, e.g. [2, 3, 4]. KK modes from the short throat begin to tunnel into the long throat. The results of the previous section (where four dimensional outer space is not expanding) indicate that the cosmological scenario with the multiple throat geometry of the inner manifold drastically differs from what was considered earlier in the literature on warped brane

string theory cosmological model.

In fact, expansion of the universe makes the multiple-throat cosmological scenario even more involved, and in some aspects different from the scenario without cosmological expansion considered of the previous Section.

During $D3\overline{D3}$ branes inflation in the short throat the Hubble parameter is $H = \sqrt{\frac{2T_3}{3M_p^2}} \sim 10^{13}\text{GeV}$. After inflation this value decreases with time as $1/t$. Thus, initially expansion of the universe is significant and would dilute the energy density of KK modes in the long throat, which are much lighter than that of the short throat ($m_{2KK} \ll m_{1KK}$). KK modes of the long throat behave as radiation, while KK modes in the short throat have matter equation of state. As a result the tunneling will happen only after the Hubble rate drops below the value equal to the tunneling rate [7], $H \sim \Gamma_{tun} \sim 10\text{GeV}$. Energy density in the long throat at this moment is

$$\epsilon_2 \sim M_p^2 H_{tun}^2 \sim M_p^2 \Gamma_{tun}^2 \sim M_p^2 m_{1KK}^2 e^{-8A_1} , \quad (5.1)$$

in our example $\epsilon_2 \sim (10^9 \text{GeV})^4$.

However, large value of the Hubble parameter generates mass gap of the KK modes $\Delta m_{KK}^2 = 2H^2$. For instance, it is known that massive gravitons in the four-dimensional de Sitter geometry have the mass gap $2H^2$ [34]. As long as H is large, KK modes of the long throat are not light, but become lighter and lighter as H decreases. As H drops below Γ_{tun} , those KK modes can be treated as radiation and we return to the same estimate (5.1) ¹².

Comparing (5.1) and (3.13) we find

$$\frac{\epsilon_2}{\epsilon_{2c}} \simeq \frac{M_p^2 m_{1KK}^2}{T_3} \frac{(e^{-A_1})^8}{(e^{-A_2})^4} \frac{1}{2e^{1+2k_c} M K} , \quad (5.2)$$

which is 10^{28} in our example. Thus tunneling KK modes from the short throat collapse into a black hole in the long throat, see Figure 1. Our universe enters the *Black Universe* (BU) phase of its cosmological evolution. Because of the horizon, dimensional reduction (1.1) is not relevant. Meanwhile, high dimensional strong gravity theory is very complicated.

Yet, the further evolution of the Black Universe phase can be understood with the magic of duality. Had the long throat cloaked with the horizon been infinitely long,

¹²Above estimate is correct provided that the problem of the angular KK modes [5] is resolved and the decay of KK modes into gravitons is suppressed compare to the tunneling time, see [7] for details. Our choice of parameters respect these conditions

it would have been holographically dual to the deconfined Klebanov-Tseytlin plasma (discussed in Section 2) in Minkowski space-time without gravity. In this case, however, the four dimensional Planck mass M_p would be infinitely large. Gluing the long throat to a compact manifold produces a finite M_p , and in the dual picture this corresponds to the coupling of the hot KT plasma to 4d gravity [35, 36, 19]. In other words, the compact manifold with strong gravity in two throats is dual to the four dimensional theory (1.3) where we will add the four-dimensional gravity R_4 . Energy density of the deconfined KT plasma will drive adiabatic expansion of the background space-time, while redshifting itself as radiation. Such an expansion continues until the plasma energy density redshifts to the critical energy density of the first order confinement transition (3.13). The Hubble scale at the beginning of radiation-dominated expansion is

$$H_{initial}^{BU} \sim \sqrt{\frac{\epsilon_2}{M_p^2}} \sim m_{1KK} e^{-4A_1}, \quad (5.3)$$

while the Hubble scale at the confinement transition is

$$H_{final}^{BU} \sim \sqrt{\frac{\epsilon_{2c}}{M_p^2}} \propto 10^{-14} H_{initial}^{BU}. \quad (5.4)$$

Since

$$\frac{H_{initial}^{BU}}{\epsilon_2^{1/4}} \sim \sqrt{\frac{m_{1KK}}{M_p}} (e^{-A_1})^2 \propto 10^{-9}, \quad \frac{H_{final}^{BU}}{\epsilon_{2c}^{1/4}} \sim 10^{-16}, \quad (5.5)$$

to an excellent approximation expanding KT plasma can be considered to be in thermal equilibrium in (almost) flat space-time during the whole period to expansion, up to the confinement phase transition.

For the Black Universe evolution dual to the KT plasma expansion (cooling) with the subsequent first order confinement phase transition, we expect that the energy density of the black hole in the long throat will dilute, according to the four-dimensional radiation dominated cosmology. In higher dimensional picture this corresponds to the recession of the horizon in the direction towards imaginary tip of the throat. As the black hole horizon energy density redshifts below the critical energy density, a first order phase transition must take place that would remove the horizon from the long throat and expose the tip of the throat¹³. The disappearance of the horizon will occur in patches in accordance with the dual picture of the first order phase transitions.

¹³A 'horizon removal' transition was discussed previously in [20]. It is not clear that our transition is similar. We comment in the conclusion how one can study such transition in the context of gauge/string duality.

Instead of the geometry with the horizon, the long throat will be filled up with the KK excitations. If there were light SM fields, say, attached to probe brane(s) around the tip of the long throat, they would be produced due to the decay these KK modes. It is at this stage that our universe is 'born', and the SM hot FRW cosmology follows. In this case one could estimate the reheating temperature of the universe in the scenario. The energy density at the phase transition is $\epsilon_{2c} \sim (100\text{GeV})^4$. Assuming that all this energy is available for the reheating of the SM, $\epsilon_{2c} \sim \frac{\pi^2}{30} g_* T_{RH}^4$ (where $g_* \propto 10^2$ is the number of the Standard Model degrees of freedom), we find a relatively low reheating temperature $T_{RH} \sim 50 \text{ GeV}$.

However, this may be irrelevant, because there is a new problem in the cosmological scenario with hierarchical throats. Indeed, it was assumed for successful phenomenology that the probe brane(s) containing SM field are located in the long throat from the very beginning. Meanwhile, emergence of the Black Brane horizon cloaks the long throat together with any probe branes located there. In a sense, the SM sector becomes screened from the theory. Later on, the geometrical phase transition with melting horizon and re-appearance of the compact throat solution is not accompanied by the re-appearance of the SM brane. The SM sector may be missing after the "Black Universe" epoch.

6 Conclusion

We presented a cosmological scenario when excited KK modes produced from the brane/antibrane annihilation in the inflationary throat tunnel to a Standard Model throat and cloak it with a Schwarzschild horizon. This suggests a model when string theory inflation is followed the 'Black Universe' epoch. This is significantly different from the models of reheating after string theory inflation previously considered in the literature. However, dual picture suggest that the 'Black Universe' epoch can be simply described by the expanding universe filled with the hot plasma composed of light particles of a (hidden sector) gauge theory 2 in deconfinement phase with $K(T)$ degrees of freedom. These originate from the rapid decay of the massive particles of another, confined (hidden sector) gauge theory 1. After the first order phase transition, the gauge theory 2 is described by the confined phase. If light SM particles are present in the theory, the corresponding particles (glueballs) decay into SM particles with relatively low reheat temperature $10 - 100\text{GeV}$.

However, for the hierarchical multiple throat scenarios we identify the problem of the SM sector disappearance: the probe brane with SM fields, which was initially placed in the long throat, will be absorbed by the horizon together with a segment of the long throat. After the end of the end “Black Universe” epoch and re-appearance of the long throat geometry, it is not clear how the probe could re-appear. Similarly, it is not clear to us what is the dual gauge theory interpretation of the screening and re-appearance of the SM sector in the theory like (1.3).

There are various directions for further study. First, one has to resolve the problem of the SM sector. One potential resolution would be to keep SM brane far enough from the tip of the long throat so that it will be not swallowed by the horizon. On the cosmology side, we expect a very drastic phenomena to occur at the epochs of the black brane ‘horizon emergence’ as well as the ‘horizon removal’ transitions.

‘Horizon removal’ phase transition is dual to the confinement/deconfinement phase transition of the long-throat dual gauge theory plasma, coupled to the 4d gravity. How one would study such a phase transition on the gravity side of the gauge/string correspondence? In [27] it was proposed how to study the gravity dual to a gauge theory in de-Sitter space-time. In our cosmological scenario the dual gauge theory plasma couples to an FRW cosmology, driven by the plasma energy density. Thus, extending ideas of [27], we should try to set-up the boundary metric to that of the appropriate FRW cosmology. The full ten-dimensional geometry should then be reconstructed requiring the nonsingularity, along the line of the gravity dual to the boost-invariant expansion of the $\mathcal{N} = 4$ supersymmetric Yang-Mills plasma [37, 38, 39, 40].

Acknowledgments

We would like to thank Neil Barnaby, Renata Kallosh, Igor Klebanov, Volodya Miransky, Rob Myers, Marco Peloso, Joe Polchinski, Sergei Prokushkin and Dam Son for valuable discussions. We would like to thank KITP Santa Barbara for hospitality, where this work was initiated during the Program ‘Nonequilibrium Dynamics in Particle Physics and Cosmology’. AB’s research at Perimeter Institute is supported in part by the Government of Canada through NSERC and by the Province of Ontario through MRI. AB gratefully acknowledges further support by an NSERC Discovery grant and support through the Early Researcher Award program by the Province of Ontario. LK was supported by NSERC and CIFAR.

References

- [1] S. Kachru, R. Kallosh, A. Linde and S. P. Trivedi, Phys. Rev. D **68**, 046005 (2003) [arXiv:hep-th/0301240].
- [2] N. Barnaby, C. P. Burgess and J. M. Cline, JCAP **0504**, 007 (2005) [arXiv:hep-th/0412040].
- [3] L. Kofman and P. Yi, Phys. Rev. D **72**, 106001 (2005) [arXiv:hep-th/0507257].
- [4] A. R. Frey, A. Mazumdar and R. C. Myers, Phys. Rev. D **73**, 026003 (2006) [arXiv:hep-th/0508139].
- [5] J. F. Dufaux, L. Kofman and M. Peloso, arXiv:0802.2958 [hep-th].
- [6] X. Chen and S. H. Tye, JCAP **0606**, 011 (2006) [arXiv:hep-th/0602136].
- [7] D. Chialva, G. Shiu and B. Underwood, JHEP **0601**, 014 (2006) [arXiv:hep-th/0508229].
- [8] B. v. Harling, A. Hebecker and T. Noguchi, JHEP **0711**, 042 (2007) [arXiv:0705.3648 [hep-th]].
- [9] S. Kachru, R. Kallosh, A. Linde, J. M. Maldacena, L. McAllister and S. P. Trivedi, JCAP **0310**, 013 (2003) [arXiv:hep-th/0308055].
- [10] I. R. Klebanov and M. J. Strassler, JHEP **0008**, 052 (2000) [arXiv:hep-th/0007191].
- [11] E. Witten, Adv. Theor. Math. Phys. **2**, 505 (1998) [arXiv:hep-th/9803131].
- [12] I. R. Klebanov and A. A. Tseytlin, Nucl. Phys. B **578**, 123 (2000) [arXiv:hep-th/0002159].
- [13] A. Buchel, Nucl. Phys. B **600**, 219 (2001) [arXiv:hep-th/0011146].
- [14] A. Buchel, C. P. Herzog, I. R. Klebanov, L. A. Pando Zayas and A. A. Tseytlin, JHEP **0104**, 033 (2001) [arXiv:hep-th/0102105].
- [15] S. S. Gubser, C. P. Herzog, I. R. Klebanov and A. A. Tseytlin, JHEP **0105**, 028 (2001) [arXiv:hep-th/0102172].

- [16] O. Aharony, A. Buchel and P. Kerner, Phys. Rev. D **76**, 086005 (2007) [arXiv:0706.1768 [hep-th]].
- [17] O. Aharony, Y. E. Antebi and M. Berkooz, Phys. Rev. D **72**, 106009 (2005) [arXiv:hep-th/0508080].
- [18] M. K. Benna, A. Dymarsky, I. R. Klebanov and A. Solovvov, arXiv:0712.4404 [hep-th].
- [19] N. Arkani-Hamed, M. Porrati and L. Randall, JHEP **0108**, 017 (2001) [arXiv:hep-th/0012148].
- [20] P. Creminelli, A. Nicolis and R. Rattazzi, JHEP **0203**, 051 (2002) [arXiv:hep-th/0107141].
- [21] I. R. Klebanov and N. A. Nekrasov, Nucl. Phys. B **574**, 263 (2000) [arXiv:hep-th/9911096].
- [22] M. Krasnitz, arXiv:hep-th/0011179.
- [23] O. Aharony, A. Buchel and A. Yarom, JHEP **0611**, 069 (2006) [arXiv:hep-th/0608209].
- [24] O. Aharony, A. Buchel and A. Yarom, Phys. Rev. D **72**, 066003 (2005) [arXiv:hep-th/0506002].
- [25] N. Seiberg, Nucl. Phys. B **435**, 129 (1995) [arXiv:hep-th/9411149].
- [26] A. Buchel and A. A. Tseytlin, Phys. Rev. D **65**, 085019 (2002) [arXiv:hep-th/0111017].
- [27] A. Buchel, Phys. Rev. D **65**, 125015 (2002) [arXiv:hep-th/0203041].
- [28] A. Buchel and R. Roiban, Phys. Lett. B **590**, 284 (2004) [arXiv:hep-th/0311154].
- [29] O. Aharony and A. Buchel, to appear.
- [30] S. B. Giddings, S. Kachru and J. Polchinski, Phys. Rev. D **66**, 106006 (2002) [arXiv:hep-th/0105097].
- [31] S. Dimopoulos, S. Kachru, N. Kaloper, A. E. Lawrence and E. Silverstein, Phys. Rev. D **64**, 121702 (2001) [arXiv:hep-th/0104239].

- [32] M. W. Choptuik, Phys. Rev. Lett. **70** (1993) 9.
- [33] T. Shiromizu, K. i. Maeda and M. Sasaki, Phys. Rev. D **62**, 024012 (2000) [arXiv:gr-qc/9910076]; S. Mukohyama, Phys. Lett. B **473**, 241 (2000) [arXiv:hep-th/9911165].
- [34] A. Higuchi, Class. Quant. Grav. **8**, 1961 (1991); **8**, 1983 (1991).
- [35] E. Witten, unpublished, H. L. Verlinde, Nucl. Phys. B **580**, 264 (2000) [arXiv:hep-th/9906182].
- [36] S. S. Gubser, Phys. Rev. D **63**, 084017 (2001) [arXiv:hep-th/9912001].
- [37] R. A. Janik and R. Peschanski, Phys. Rev. D **73**, 045013 (2006) [arXiv:hep-th/0512162].
- [38] R. A. Janik, Phys. Rev. Lett. **98**, 022302 (2007) [arXiv:hep-th/0610144].
- [39] M. P. Heller and R. A. Janik, Phys. Rev. D **76** (2007) 025027 [arXiv:hep-th/0703243].
- [40] P. Benincasa, A. Buchel, M. P. Heller and R. A. Janik, Phys. Rev. D **77**, 046006 (2008) [arXiv:0712.2025 [hep-th]].